

Dark Matter from Split Seesaw

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Abstract

The seesaw mechanism in models with extra dimensions is shown to be generically consistent with a broad range of Majorana masses. The resulting democracy of scales implies that the seesaw mechanism can naturally explain the smallness of neutrino masses for an arbitrarily small right-handed neutrino mass. If the scales of the seesaw parameters are split, with two right-handed neutrinos at a high scale and one at a keV scale, one can explain the matter-antimatter asymmetry of the universe, as well as dark matter. The dark matter candidate, a sterile right-handed neutrino with mass of several keV, can account for the observed pulsar velocities and for the recent data from *Chandra* X-ray Observatory, which suggest the existence of a 5 keV sterile right-handed neutrino.

I. INTRODUCTION

The possible existence of extra dimensions has changed the way one thinks about the hierarchy of scales in particle physics [1, 2]. If the light Standard Model (SM) degrees of freedom are localized on a (3+1)-dimensional brane embedded in a higher-dimensional spacetime, while any other degrees of freedom reside on other branes or propagate in the bulk, the perceived difference between the high scale and the low scale may arise from the suppression that is exponential in the size of the extra dimension. It is a generic phenomenon in this class of models that the interactions of fields localized on different branes can be exponentially suppressed by the distance between the branes [3]. In the low-energy effective theory, the smallness of the couplings may seem completely unnatural, because the symmetry group of the (3+1)-dimensional effective theory need not increase in the limit where the size of the extra dimension becomes large.

Neutrino masses and the possibility of leptogenesis offer an intriguing connection to the high-scale physics. In (3+1) dimensions, the naturalness arguments suggest that the $SU(2)$ singlet fermions should have masses of the order of the high scale, in contrast with the non-singlet fermions, whose masses are protected by the chiral gauge $SU(2)$ symmetry. However, if extra dimensions exist, one must rethink the naturalness arguments. A fermion localized on a distant brane may appear as weakly interacting light particle in the low-energy effective four-dimensional Lagrangian.

Let us review the natural values of parameters in the modern version of the Standard Model, by which we mean the original Standard Model of Glashow, Salam, and Weinberg, supplied with three right-handed neutrinos generating the observed neutrino masses via the seesaw mechanism [4]. As mentioned above, and as we will show in detail below, both heavy and light right-handed neutrinos can be equally natural if extra dimensions exist, assuming that the small-mass fermions in (3+1) dimensions arise from the heavy particles on a remote brane. The simultaneous existence of light and heavy right-handed neutrinos allows one to explain both the baryon asymmetry of the universe and dark matter in the framework of the simplest model consistent with the data, – the modern variant of the Standard Model. The dark matter candidate in this model is a keV sterile neutrino, the possible existence of which is supported by some astrophysical data [5–8].

The only way to generate the baryon asymmetry of the universe in this minimal scenario

is via leptogenesis [9], because the lack of CP violation and the weakness of the electroweak phase transition eliminate any alternative scenarios. Only two right-handed neutrinos are necessary for leptogenesis [10]. However, having three right-handed neutrinos, one for each generation of fermions, is arguably more appealing; this also allows one to embed the Standard Model into SO(10) Grand Unified Theory or any theory that has a gauge $U(1)_{B-L}$ symmetry. What is the natural value for the Majorana mass of the third right-handed field? In the presence of extra dimensions, the high-scale values are not necessarily favored by any naturalness arguments. However, if this Majorana mass is of the order of keV, one can explain dark matter, as well as the long-standing puzzle of the pulsar kicks [5, 6]. The same particles can play an important role in the formation of the first stars [7] and other astrophysical phenomena [11]. Furthermore, at least two recent astrophysical observations suggest the existence of a sterile neutrino either with 5 keV mass [8] or 17 keV mass [12].

By splitting the scales of the heavy and light right-handed neutrinos, which is naturally realized in the *split seesaw* model described below, we achieve an elegant and simple description of all known experimental data. We note that models with three right-handed neutrinos below the electroweak scale have been proposed [13, 14], and in one of them, dubbed ν MSM [14], one can explain the baryon asymmetry by demanding a high degree of degeneracy between two Majorana masses, both of the order of several GeV. The high degree of degeneracy is required to amplify the effects of CP violation in a leptogenesis scenario that involves oscillations [15]. In contrast, we employ the standard leptogenesis using decays of the heavy right-handed neutrinos, and no mass degeneracy is required. We emphasize, however, that our model with an extra dimension is not limited to the conventional leptogenesis scenario, and it can explain a very broad range of scales. Indeed, our model assures that the beauty of the seesaw formula, which explains the smallness of the neutrino masses by relating them to the ratio of the weak scale and the GUT scale, is preserved for practically arbitrary choice of the right-handed neutrino masses.

II. SPLIT SEESAW MECHANISM

We begin by reviewing the seesaw mechanism in the 4D theory. The relevant terms in the Lagrangian are given by

$$\mathcal{L} = i\bar{N}_i\gamma^\mu\partial_\mu N_i + \left(\lambda_{i\alpha}\bar{N}_i L_\alpha \phi - \frac{1}{2}M_{Ri}\bar{N}_i^c N_i + \text{h.c.} \right), \quad (1)$$

where N_i , L_α and ϕ are the right-handed neutrino, lepton doublet and Higgs boson, respectively, i denotes the generation of the right-handed neutrino, and α runs over the lepton flavor, e , μ and τ . Integrating out the massive right-handed neutrinos yields the seesaw formula for the light neutrino mass:

$$(m_\nu)_{\alpha\beta} = \sum_i \lambda_{i\alpha}\lambda_{i\beta} \frac{\langle\phi^0\rangle^2}{M_{Ri}}. \quad (2)$$

The atmospheric [16] and solar [17, 18] neutrino oscillation experiments have provided firm evidence that at least two neutrinos have small but non-zero masses, and the mass splittings are given by $\Delta m_{\text{atm}}^2 \simeq 2 \times 10^{-3} \text{eV}^2$ and $\Delta m_{\odot}^2 \simeq 8 \times 10^{-5} \text{eV}^2$. The seesaw mechanism then suggests that a typical mass scale of the right-handed neutrinos should be $\sim 10^{15} \text{GeV}$, close to the GUT scale, for $\lambda_{i\alpha} \sim 1$.

To explain the neutrino oscillation data, one must introduce more than one right-handed neutrino. It was shown in Ref. [10] that two right-handed neutrinos, N_2 and N_3 , would suffice for this purpose. Moreover, it is possible to generate the cosmological baryon asymmetry via leptogenesis with the two right-handed neutrinos. Thus, the addition of the two right-handed neutrinos appears to be the most economical extension of the SM.

However there are two issues in the above model, which, as we will see below, can be addressed simultaneously and naturally in a theory with an extra dimension. One issue is that at least the lightest of $N_{2,3}$ must have a mass $O(10^{11-12}) \text{GeV}$, several orders of magnitude below the GUT scale, for the leptogenesis to work [19, 20]. Therefore, to explain the neutrino masses, one must suppress both the right-handed Majorana mass and the associated Yukawa couplings. This is not a severe fine-tuning, and it can be rectified by the Froggatt-Nielsen mechanism [21], but the introduction of a new symmetry and the breaking of this symmetry at an appropriate scale, for this purpose, appear somewhat *ad hoc*. The other issue is that, if one gauges the $U(1)_{B-L}$ symmetry, three right-handed neutrinos are required to achieve the anomaly cancellation. Then the question arises: what is the role of

the third right-handed neutrino, N_1 ? In one limiting case, N_1 may have a Planck-scale mass, and this particle would play no role in the low energy physics. In particular, it is clear from the seesaw formula that such a heavy N_1 does not contribute to the light neutrino mass, m_ν . In the other limit, N_1 may be very light. An interesting possibility arises in this case; if the mass is at a keV scale, then N_1 , which is long-lived on cosmological time scales, is a viable dark-matter candidate [22]. In addition, the same particle would be produced in a supernova explosion, and it would be emitted with an anisotropy of a few per cent, which is sufficient to explain the observed velocities of pulsars [6]. However, one could ask whether such an extremely light mass may upset the seesaw mechanism, and it is also unclear how such a small mass may arise naturally. We will show that, in a theory with an extra dimension, a split mass spectrum arises naturally, without disrupting the seesaw mechanism. The *split seesaw* mechanism can explain both mild and large mass hierarchy, i.e., why $M_{2,3} < M_{\text{GUT}}$ and why $M_1 \ll M_{2,3}$.

Let us consider a 5D theory compactified on S^1/Z_2 with coordinate $y \in [0, \ell]$. One of the boundaries at $y = 0$ is identified with the SM brane, where the SM degrees of freedom reside, while the other at $y = \ell$ is a hidden brane. The size of the extra dimension ℓ and the 5D fundamental scale M are related to the 4D reduced Planck scale as

$$M_P^2 = M^3 \ell. \quad (3)$$

For the consistency of the theory, the compactification scale $M_c \equiv 1/\ell$ must be smaller than M . As reference values we take $M \sim 5 \times 10^{17}$ GeV, $M_c \sim 10^{16}$ GeV and the $B - L$ breaking scale $v_{B-L} \sim 10^{15}$ GeV in the following.

We introduce a Dirac spinor field in 5D, $\Psi = (\chi_\alpha, \bar{\psi}^{\dot{\alpha}})^T$, with a bulk mass m :

$$S = \int d^4x dy M (i\bar{\Psi}\Gamma^A\partial_A\Psi + m\bar{\Psi}\Psi), \quad (4)$$

where A runs over 0, 1, 2, 3, 5, and the 5D gamma matrices Γ^A are defined by

$$\Gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix}, \quad \Gamma^5 = -i \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (5)$$

We have inserted the mass scale M in Eq. (4) so that the mass dimension of Ψ is 3/2 as in the 4D case. The zero mode of Ψ should satisfy

$$(i\Gamma^5\partial_5 + m)\Psi^{(0)} = 0. \quad (6)$$

The bulk profile of the zero mode is therefore given by $\exp(\mp my)$ for χ and $\bar{\psi}$ [23, 24].

A 4D chiral fermion can be obtained from the 5D Dirac fermion by orbifolding. If we assign a Z_2 parity -1 and $+1$ to χ and ψ , respectively, we can see that only ψ has a zero mode with an exponential profile in the bulk. (For consistency we have assigned a negative Z_2 parity to the bulk mass m .) The zero mode of $\Psi_R = (0, \bar{\psi})^T$ can be expressed in terms of the canonically normalized (right-handed) fermion $\psi_R^{(4D)}$ in the 4D theory as

$$\Psi_R^{(0)}(y, x) = \sqrt{\frac{2m}{e^{2m\ell} - 1}} \frac{1}{\sqrt{M}} e^{my} \psi_R^{(4D)}(x). \quad (7)$$

As one can see from the y -dependence of the wave function, the zero mode peaks at the hidden brane at $y = \ell$ for positive m . In particular, in the limit of $m\ell \gg 1$, the overlap of the zero mode with particles on the SM brane becomes exponentially suppressed.

We would like to promote the right-handed neutrino N_i to a bulk field Ψ_i in 5D [3]. For definiteness, we identify the zero modes of Ψ_{iR} with the right-handed neutrinos in the 4D theory. Then the zero modes have an exponential profile in the y direction $\sim \exp(my)$. We also introduce a $U(1)_{B-L}$ gauge field in the bulk and a scalar field Φ with a $B - L$ charge -2 on the SM brane. Assuming that the Φ develops a VEV $v_{B-L} \approx 10^{15} \text{ GeV}$, the zero mode of the $U(1)_{B-L}$ gauge boson receives mass via the usual Higgs mechanism, and there is no additional zero mode. The VEV of Φ also gives rise to heavy Majorana masses for the right-handed neutrinos. After integrating out the heavy Φ and the $U(1)_{B-L}$ gauge boson, one obtains the Lagrangian for the zero modes:

$$S = \int d^4x dy \left\{ M \left(i \bar{\Psi}_{iR}^{(0)} \Gamma^A \partial_A \Psi_{iR}^{(0)} + m_i \bar{\Psi}_{iR}^{(0)} \Psi_{iR}^{(0)} \right) + \delta(y) \left(\frac{\kappa_i}{2} v_{B-L} \bar{\Psi}_{iR}^{(0)c} \Psi_{iR}^{(0)} + \tilde{\lambda}_{i\alpha} \bar{\Psi}_{iR}^{(0)} L_\alpha \phi + \text{h.c.} \right) \right\}, \quad (8)$$

where κ_i and $\tilde{\lambda}_{i\alpha}$ are numerical constants of order unity, and we introduced the lepton and Higgs doublets on the SM brane at $y = 0$.^{#1} Using the relation (7), we obtain the effective 4D mass and Yukawa couplings in (1) as [26],

$$M_{Ri} = \kappa_i v_{B-L} \frac{2m_i}{M(e^{2m_i\ell} - 1)}, \quad (9)$$

$$\lambda_{i\alpha} = \frac{\tilde{\lambda}_{i\alpha}}{\sqrt{M}} \sqrt{\frac{2m_i}{e^{2m_i\ell} - 1}} = \tilde{\lambda}_{i\alpha} \sqrt{\frac{M_{Ri}}{\kappa_i v_{B-L}}}. \quad (10)$$

^{#1} Since all the right-handed neutrinos are in the bulk, the $U(1)_{B-L}$ becomes anomalous on the SM brane. However, one can cancel the anomaly by introducing Chern-Simons terms with an appropriate coefficient in the bulk [25].

As expected, the right-handed neutrino mass and the Yukawa couplings are suppressed by an exponential factor. In the extreme limit of $m_i \ell \gg 1$, both M_R and λ are extremely suppressed. What is remarkable is that the exponential factors cancel in the seesaw formula for the light neutrino masses (2) because the Yukawa coupling squared is suppressed by the same factor as the right-handed mass [3]. Thus, the neutrino masses are given by the familiar seesaw relation:

$$(m_\nu)_{\alpha\beta} = \left(\sum_i \frac{1}{\kappa_i} \tilde{\lambda}_{i\alpha} \tilde{\lambda}_{i\beta} \right) \frac{\langle \phi^0 \rangle^2}{v_{B-L}}, \quad (11)$$

where the quantity in the parenthesis is of order unity. We can clearly see that no small parameters such as exponentially suppressed mass or coupling appear in the seesaw formula; the typical neutrino mass scale is given by the ratio of the square of the weak scale to the $B - L$ breaking scale, showing that the seesaw mechanism is robust against splitting the right-handed neutrino mass spectrum. So we can easily realize a split mass spectrum by choosing appropriate (not extremely large or small) values of m_i , without spoiling the seesaw mechanism. For instance, we obtain $M_{R2} \sim 10^{12}(10^{11}) \text{ GeV}$ and $M_{R1} \sim \text{keV}$ for $m_2 \simeq 2.3(3.6)\ell^{-1}$ and for $m_1 \simeq 24\ell^{-1}$, respectively, where we set $\kappa_i = 1$ and used the reference values of v_{B-L} , ℓ and M .^{#2}

III. COSMOLOGY

In the context of sterile neutrino dark matter, it is customary to parametrize it in terms of the effective mixing angle θ and the mass m_s . They are given by

$$m_s = M_{R1}, \quad (12)$$

$$\begin{aligned} \theta^2 &\simeq \frac{\sum_\alpha |\lambda_{1\alpha}|^2}{M_{R1}^2} \langle \phi^0 \rangle^2 \\ &\simeq 3 \times 10^{-9} \left(\frac{\kappa_1^{-1} \sum_\alpha |\tilde{\lambda}_{1\alpha}|^2}{10^{-4}} \right) \left(\frac{10^{15} \text{ GeV}}{v_{B-L}} \right) \left(\frac{1 \text{ keV}}{m_s} \right), \end{aligned} \quad (13)$$

$$m_\nu \simeq \theta^2 m_s, \quad (14)$$

where we have assumed $\theta \ll 1$. If N_1 is the dark-matter particle, the X-ray constraints imply $\theta^2 \lesssim 10^{-5}(m_s/\text{keV})^{-4}$, while the small-scale structure constraints imply $m_s > 2 \text{ keV}$

^{#2} It is also possible to realize $M_{R,2} \simeq M_{R3} \sim \text{GeV}$ for $m_2 \simeq m_3 \simeq 17\ell^{-1}$, as considered in the νMSM [14].

for the resonant production [27] (see, e.g., Refs. [5, 28] for review). For the above range of parameters the contribution of N_1 to the light neutrino mass is negligible [29]. The prospects of direct detection of sterile dark matter are also stymied by the smallness of the mixing angle [30]. In contrast, the X-ray instruments can search for decays of sterile right-handed neutrinos, which have the half-life much longer than the age of the universe. Here the smallness of $\sin^2 \theta$ is compensated by the huge number of dark-matter particles. In particular, if the 2.5 keV spectral feature in the recent *Chandra* X-ray Observatory data [8] is explained by the decay of a 5 keV relic sterile right-handed neutrino, the inferred parameters $m_s \simeq 5 \text{ keV}$ and $\theta = (0.2 - 1.4) \times 10^{-9}$ [8] imply that such a sterile neutrino does not contribute to the neutrino mass.

To account for the observed dark matter density, sterile right-handed neutrinos must have the correct abundance. A generic way to produce relic sterile neutrinos is through non-resonant (NR) oscillations [22]. For the mixing angle $\sin^2 \theta \approx 10^{-9}$ and the mass $m_s \approx 5 \text{ keV}$, the resulting abundance of sterile neutrinos is in the right range [22, 31–33]. Assuming that the sterile neutrinos produced by the NR oscillations account for the total dark matter density, there is an upper limit $m_s \lesssim 5 \text{ keV}$ from the X-ray observations [34]. On the other hand, the lower limits $m_s \gtrsim 8 \text{ keV}$ [35] and $m_s \gtrsim 11 \text{ keV}$ [36] based on the Ly- α observations and on dwarf spheroidals, respectively, appear to rule out such dark matter produced by NR oscillations. Thus, if the X-ray line observed by the *Chandra* X-ray Observatory is due to the decay of dark matter sterile neutrinos, they must be produced by some other mechanism. Such mechanisms, producing colder sterile neutrinos, consistent with all the constraints, have been studied in the context of other models [37]. We will show that, in the split seesaw scenario, the sterile neutrinos can be produced with the same or greater cool-down factor as those produced at the electroweak scale [37]. Thus, our dark-matter candidate presents no conflict with the Lyman- α and other small-scale structure bounds.

In our set-up, the sterile neutrino has a $U(1)_{B-L}$ gauge interaction, which may become important in the early Universe. In particular, the reheat temperature T_R must be greater than M_{R2} or M_{R3} for the leptogenesis to work. According to Ref. [20], a successful leptogenesis is possible for $T_R \gtrsim 10^{11} \text{ GeV}$, unless M_{R2} and M_{R3} are extremely degenerate. For such a high reheat temperature, the main production process is pair production of N_1 from the SM fermions in plasma through the s-channel exchange of the B-L gauge boson. The production is most efficient at reheating. The number to entropy ratio of the sterile neutrino

produced by this mechanism can be roughly estimated as

$$\begin{aligned}
Y_{N1} &\equiv \frac{n_{N1}}{s} \sim \left. \frac{\langle \sigma v \rangle n_f^2 / H}{\frac{2\pi^2}{45} g_* T^3} \right|_{T=T_R} \\
&\sim 10^{-4} \left(\frac{g_*}{10^2} \right)^{\frac{3}{2}} \left(\frac{v_{B-L}}{10^{15} \text{ GeV}} \right)^{-4} \left(\frac{T_R}{5 \times 10^{13} \text{ GeV}} \right)^3,
\end{aligned} \tag{15}$$

where H is the Hubble parameter, g_* counts the relativistic degrees of freedom at the reheating, $\langle \sigma v \rangle \sim T^2/v_{B-L}^4$ is the production cross section, $n_f \sim T^3$ is the number density of the SM fermions in plasma, and the first equality is evaluated at the reheating. The numerical solution of the Boltzmann equation gives a consistent result [38]. The dark matter abundance is related to the mass density to entropy ratio as

$$\Omega_{\text{dark}} = 0.2 \times \left(\frac{m_s}{5 \text{ keV}} \right) \left(\frac{Y_{N1}}{0.7 \times 10^{-4}} \right). \tag{16}$$

Thus, the reheat temperature as large as 10^{13} GeV is needed to account for all the dark matter density by this production process. The thermal leptogenesis works with such a high temperature. Also since the sterile right-handed neutrinos are out of equilibrium since their production at high temperature, when all the Standard model degrees of freedom are thermally excited, the average momentum of such dark-matter particles is red-shifted by a factor of a cubic root of the ratio of the degrees of freedom, which is about 3.5 [37].

Another scenario for producing dark matter with the correct abundance (but with different kinetic properties) places no constraint on the reheat temperature and predicts an even colder dark matter, as long as reheating restores the gauge $U(1)_{B-L}$ symmetry. We assume that the $U(1)_{B-L}$ symmetry is broken spontaneously at 10^{15} GeV by the VEV of the Higgs boson Φ . Since the right-handed neutrinos are coupled to the $U(1)_{B-L}$ gauge boson with the universal gauge coupling g_0 , they are produced at high temperature and they reach the equilibrium density $n_{\text{eq}}(T)$ with a distribution function $f_{\text{eq}}(E, T) = 1/(\exp(E/T) + 1)$ above the transition temperature T_t . However, after the phase transition, the $U(1)_{B-L}$ gauge boson becomes massive and the production of N_i is suppressed by the gauge boson mass (just as in the production mechanism discussed above). Meanwhile, the rest of the Standard Model particles are produced in the true vacuum. The entropy produced in this phase transition can dilute the number density of dark-matter sterile neutrinos from its equilibrium value $Y_{N1}^{\text{eq}} = 45\zeta(3)2/(2\pi^4 g_*) \sim 10^{-2}$ to what is required to explain dark matter, see eq. (16). In this case, the reheat temperature is not constrained from above, but the final temperature

after the phase transition T_f must satisfy the same constraint as the reheat temperature in the previous scenario.

If the Higgs potential at zero temperature is

$$V(\phi, 0) = -\frac{\mu^2}{2}\phi^2 + \frac{\lambda}{4}\phi^4, \quad (17)$$

where we fix the vev of the Higgs to be v_{B-L} , i.e.,

$$\frac{\mu}{\sqrt{\lambda}} = v_{B-L} \approx 10^{15} \text{ GeV}. \quad (18)$$

The temperature-dependent effective potential [39] is

$$V(\phi, T) \approx V(\phi, 0) + \frac{3g^2 + 4\lambda}{24}T^2\phi^2 - \frac{3g^3 + g\lambda + 3\lambda^{3/2}}{24\pi}T\phi^3, \quad (19)$$

where $g = 2g_0$ is the effective gauge $(B-L)$ coupling of the Higgs boson, which has $(B-L) = -2$. The phase transition from the $U(1)$ symmetric vacuum to the broken-symmetry vacuum takes place when the tunneling rate [39, 40] per Hubble volume per Hubble time is of order one:

$$T_t^4 \left(\frac{S_3}{2\pi T} \right)^{\frac{3}{2}} e^{-S_3/T_t} \times H_t^{-4} \sim 1, \quad (20)$$

where

$$H_t = \left(\frac{\pi^2 g_*}{90} \right)^{1/2} \frac{T_t^2}{M_P}. \quad (21)$$

Here S_3 is the action of the three-dimensional bounce, and in the thick-wall limit, it is approximately given by [39]

$$\frac{S_3}{T} \simeq \frac{44 (\gamma(T^2 - T_c^2))^{3/2}}{\alpha^2 T^3}, \quad (22)$$

with

$$\alpha \equiv \frac{3g^3 + g\lambda + 3\lambda^{3/2}}{8\pi}, \quad \gamma \equiv \frac{3g^2 + 4\lambda}{12}, \quad T_c \equiv \frac{\mu}{\sqrt{\gamma}}.$$

The correct abundance of dark matter is attained if the ratio of entropies before and after the phase transition is given by

$$\frac{S_f}{S_t} \sim \frac{T_f^3}{T_t^3} \sim 10^2, \quad \text{or} \quad T_f \sim 5T_t. \quad (23)$$

To avoid overproduction of sterile neutrino after the phase transition, the final temperature should be smaller than the reheat temperature estimated in the previous scenario,

$$T_f < T_{\max} = 5 \times 10^{13} \text{ GeV} \left(\frac{v_{B-L}}{10^{15} \text{ GeV}} \right)^{4/3}. \quad (24)$$

By energy conservation,

$$T_f \sim (\epsilon/g_*)^{1/4} \approx 0.3\epsilon^{1/4},$$

where $\epsilon = -V(v_{B-L}, 0)$. Note here that the latent heat comes from the energy difference between the symmetric phase $\phi = 0$ and the broken phase $\phi = v_{B-L}$, which is estimated by using the zero-temperature potential.

While there is a significant freedom in the choice of parameters, one can verify that, for example, the choice of $g \equiv 2g_0 = 0.3$, $\mu = 10^{11}$ GeV satisfies the constraints (23,24) for $T_t \sim 7 \times 10^{11}$ GeV and $T_f \approx 5T_t \simeq 3 \times 10^{12}$ GeV. For an entropy dilution of $O(100)$ to occur, a relatively light mass μ is required, which leads to some degree of fine-tuning of the parameters in the scalar potential. This fine-tuning may be ameliorated by extending the Higgs sector, so as to increase the strength of the first-order phase transition.

While the X-ray constraints depend only on the mass, mixing, and abundance of sterile neutrinos, the clustering properties of dark matter, probed by Lyman- α bounds [35] and the observations of dwarf spheroidal galaxies [36, 41], are sensitive to the momentum distribution of dark-matter particles. In our first scenario, the limited preheating produces a non-thermal distribution of dark matter particles, which is further cooled down and redshifted in the subsequent history of the universe. In our second scenario, in addition to this redshifting, the B-L breaking phase transition redshifts the momenta of dark-matter particles by an additional factor $T_f/T_t \approx 5$. Thus, the two scenarios produce dark matter with different free-streaming lengths, each of which is small enough to satisfy the present bounds.

IV. DISCUSSION

We have shown that the seesaw formula is robust with respect to splitting the right-handed neutrino mass spectrum in a model with an extra dimension. The Yukawa couplings and the Majorana masses are suppressed in such a way that the usual seesaw relation is preserved. This remarkable feature allows the use of the seesaw mechanism for two heavy right-handed neutrinos and one keV sterile right-handed neutrino in a model that does not require any unnatural fine-tuning of parameters. The resulting *split seesaw* model explains the origin of the baryon asymmetry of the universe, dark matter, and the smallness of the active neutrino masses in a natural and holistic manner. Furthermore, the light sterile neutrino provides a simple explanation for the observed velocities of pulsars, due

to anisotropic emission of such particles from a supernova. Finally, the model offers an explanation as to why there are three generations of fermions in nature: two right-handed neutrinos are needed for leptogenesis, and the third one plays the role of dark matter. X-ray astronomy offers an exciting opportunity to discover this dark matter candidate, and indeed the recent data from *Chandra* X-ray Observatory show a spectral feature consistent with this form of dark matter. If this result is confirmed, the detailed analysis of the structure on sub-galactic scales can help distinguish between the two production mechanisms we have discussed.

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